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PROPAGATION OF HIGH ENERGY SOLAR PROTONS IN THE INTERPLANETARY MAGNETIC FIELD

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SUMMARY

In order to describe the propagation of high energy solar protons in the interplanetary medium, the kinetic equation is applied in the Fokker-Planck approximation.

Analysis is performed of the solution of this equation so as to ascertain the conditions of observation of strong anisotropy in the angular distribution of relativistic protons generated during solar flares.

It is obtained from the comparison of theoretical results with the observation data on the flare of 4 May 1960 that the diffusion coefficient D increases with the distance from the Sun approximately as r^2. Near the terrestrial orbit $D \approx 5 \cdot 10^{22}$ cm^2/sec.

* * *

As is well known, certain flares indicate a strong anisotropy in the angular distribution of high energy solar protons in the initial stage of the flare. This means that protons, reaching the Earth in that period, undergo no substantial deviations. In this case the diffusive approximation is inapplicable and the use of a more general equation is required.

Measurements of interplanetary magnetic field on Mariner-2 show that there exists near the terrestrial orbit a magnetic field with a mean quadratic intensity of $3 \cdot 10^{-5}$ gauss and irregularity dimension from 10^{11} cm and less. On such an irregularity a relativistic proton with energy

* O RASPROSTRANENII VYSOKOENERGETICHESKIKH SOLNECHNYKH PROTONOV V MEZH — PLANENTNOM MAGNTNOM POLE

$3 \cdot 10^9$ ev scatters over an angle $< 20^\circ$. Consequently, in order to describe the propagation of high energy protons in interplanetary medium, one may apply the kinetic equation in the Fokker-Planck approximation.

The present work is devoted to its solution when applied to a solar flare; at the same time, principal attention is given the ascertaining of the conditions at which observation is possible of strong anisotropy in the angular distribution of solar protons.

BASIC EQUATION.- We shall consider that the interplanetary magnetic field consists of two components

$$\mathbf{B} = \mathbf{B}_0 + \mathbf{B}_1, \tag{1}$$

where $\mathbf{B}_0$ is the regular magnetic field that approximates large-scale fields of which the dimensions are much greater than the Larmor radius of the particle); $\mathbf{B}_1$ is a uniform random field, which will be given by the moments

$$\langle \mathbf{B}_1 \rangle = 0, \quad \langle B_{1i}(\mathbf{r}_1) \mathbf{B}_{1i}(\mathbf{r}_2) \rangle = \beta^2 \exp\left(- \frac{|\mathbf{r}_1 - \mathbf{r}_2|^2}{a^2} \right), \tag{2}$$

where $B_{1r}, B_{1\eta}, B_{1\zeta}$ are the spatial vector components in the spherical system of coordinates, $i = r, \eta, \zeta$. Angular brackets indicate the spatial averaging. The magnetic field is considered stationary.

We shall denote by $\mathbf{k}$ the unitary vector in the direction of particle motion $(k = \mathbf{v}/|\mathbf{v}|)$. We have for the accretion $\Delta \mathbf{k}$ of the vector $\mathbf{k}$ in the time interval Δt

$$\Delta \mathbf{k} = \frac{e}{mc} \int_0^{\Delta t} [\mathbf{k}' \times \mathbf{B}']\, dt \approx \frac{e}{mc} \left[\mathbf{k} \times \int_0^{\Delta t} \mathbf{B}' dt' \right]. \tag{3}$$

Here m is the total mass of the particle. Δt must be so chosen that the particle pass through many irregularities and yet with a small deviation in the direction. Effecting combinations from (3), and averaging with the utilization of (2), we shall have

$$\langle \Delta k_\varphi \rangle = \frac{eB_0}{mc} \Delta t = \omega \Delta t, \quad \langle \Delta k_\theta \rangle = \frac{e^2 \beta^2}{m^2 c^2} \frac{2\sqrt{\pi} a}{v} \operatorname{ctg} \theta \Delta t = q \operatorname{ctg} \theta \Delta t,$$

$$\langle (\Delta k_\theta)^2 \rangle = 2q \Delta t, \quad \langle (\Delta k_\varphi)^2 \rangle = \frac{2q}{\sin^2 \theta} \Delta t, \tag{4}$$

where θ and φ are the angular coordinates; θ is counted from the radius; φ is the azimuthal angle.

Utilizing these expressions for the moments, and estimating the higher order moments to be small, we shall obtain from the Markov equation, linking the value of the distribution function $N(\mathbf{r}, \theta, \varphi, t)$ at times t and t + Δt, the kinetic equation in the Fokker-Planck approximation

$$\frac{\partial N}{\partial t} + v\,\mathbf{k}\,\mathrm{grad}\,N - \frac{v\sin\theta}{r}\frac{\partial N}{\partial\theta} - \frac{v\sin\theta\sin\varphi}{r}\operatorname{ctg}\eta\frac{\partial N}{\partial\varphi} = \\ = \frac{q}{\sin\theta}\frac{\partial}{\partial\theta}\sin\theta\frac{\partial N}{\partial\theta} + \frac{q}{\sin^2\theta}\frac{\partial^2 N}{\partial\varphi^2} + \omega\frac{\partial N}{\partial\varphi}. \tag{5}$$

($N(\mathbf{r}, \theta, \varphi, t)\,d\mathbf{r}\,dk$ is the number of particles in the volume dr, having velocity directions confined in the solid angle dk).

The energy of the particle does not vary, for the field is considered stationary. The detailed development of the result is available in ref. [1].

SPHERICALLY-SYMMETRICAL FLARE.- Let us consider at the outset a spherically-symmetrical point flare in an infinite uniform medium The kinetic equation and the initial condition then have the form :

$$\frac{\partial N}{\partial t} + v\cos\theta\frac{\partial N}{\partial r} - v\frac{\sin\theta}{r}\frac{\partial N}{\partial\theta} = q\frac{1}{\sin\theta}\frac{\partial}{\partial\theta}\sin\theta\frac{\partial N}{\partial\theta}, \tag{6}$$

$$N|_{t=0} = \frac{\overline{N}}{16\pi^2 r^2}\delta(r). \tag{7}$$

Moreover, N (r, θ, t) must be a periodical and symmetrical function of θ

$$N(\theta) = N(\theta + 2\pi), \quad N(\theta) = N(-\theta). \tag{8}$$

Expanding N by Legendre polynomials

$$N = \sum_{k=0}^{\infty} a_k(r, t) P_k(\cos\theta), \tag{9}$$

we shall obtain for the coefficients a_k the chain of equations

$$\frac{\partial a_k}{\partial t} + \frac{vk}{2k-1}\frac{\partial a_{k-1}}{\partial r} + \frac{v(k+1)}{2k+3}\frac{\partial a_{k+1}}{\partial r} - \\ - \frac{v}{r}\frac{k(k-1)}{2k-1}a_{k-1} + \frac{v}{r}\frac{(k+1)(k+2)}{2k+3}a_{k+1} = -qk(k+1)a_k. \tag{10}$$

Let us write the first two equations

$$\frac{\partial a_0}{\partial t}+\frac{v}{3}\frac{\partial a_1}{\partial r}+\frac{2}{3}v\frac{a_1}{r}=0, \quad \frac{\partial a_1}{\partial t}+v\frac{\partial a_0}{\partial r}+\frac{3}{5}v\frac{\partial a_2}{\partial r}+\frac{6}{5}v\frac{a_2}{r}=-2qa_1. \tag{11}$$

The diffusive approximation is obtained in the case, when instead of the second equation of the system (11), we would write

$$2qa_1=-v\frac{\partial a_0}{\partial r}. \tag{12}$$

Note that $4\pi a_0$ is the spatial density of particles and $va_1/3$ is the flux. From the first equation (11) and from (12) we may obtain the diffusion equation with the diffusion coefficient

$$D = v^2/6q.$$

However, the equality (12) will be valid if the processes are sufficiently slow and sufficiently isotropic. A simple solution may be obtained in the case, whereby rejected in the system (11) are only the terms containing a_2, leaving the term $\partial a_1/\partial t$, that is, in the assumption that the processes are sufficiently isotropic. Thus we have the system

$$\frac{\partial a_0}{\partial t}+\frac{v}{3}\frac{\partial a_1}{\partial r}+\frac{2}{3}v\frac{a_1}{r}=0, \quad \frac{\partial a_1}{\partial t}+v\frac{\partial a_0}{\partial r}+2qa_1=0, \tag{13}$$

$$a_0|_{t=0}=\frac{\bar{N}}{16\pi^2r^2}\delta(r), \qquad a_1|_{t=0}=0. \tag{14}$$

Let us now perform the Carson-Laplace transformation

$$\bar{a}_k=z\int_0^\infty a_k e^{-zt}\,dt. \tag{15}$$

As a result we shall obtain

$$z\bar{a}_0+\frac{v}{3}\frac{1}{r^2}\frac{\partial}{\partial r}r^2\frac{\partial\bar{a}_1}{\partial r}=0, \quad (z+2q)\bar{a}_1+\frac{\partial\bar{a}_0}{\partial r}=0. \tag{16}$$

Limited at $r\to\infty$, the solution of (16) will be

$$\bar{a}_0=\frac{c}{r}\exp\left(-\frac{r}{v}\sqrt{3z(z+2q)}\right).$$

The constant c is determined from the condition

$$\int_0^\infty \bar{a}_0 r^2 dr = \frac{\bar{N}}{16\pi^2}. \tag{17}$$

Thus

$$\bar{a}_0 = \frac{3z(z+2a)}{rv^2} \exp\left[-\frac{r}{v}\sqrt{3z(z+2q)}\right]. \tag{18}$$

Passing from the image to the original, we obtain

$$a_0 = \frac{3\sqrt{3}\bar{N}}{16\pi^2 r^2}\delta(vt-\sqrt{3}r)e^{-qt} + \\ + \frac{\bar{N}q^3}{16\pi^2 v^3}\frac{3\sqrt{3}}{x}e^{-qt}\left[I_1(x) + \frac{qt}{x}I_0(x) - \frac{2qt}{x^2}I_1(x)\right]. \tag{19}$$

Analogously obtained is the expressions for a_1

$$a_1 = \frac{\bar{N}q^3}{16\pi^2 v^3}\frac{3\sqrt{3}rq}{vx^2}e^{-qt}\left[I_0(x) - \frac{2}{x}I_1(x)\right]. \tag{20}$$

Here I_1 is a Bessel function from the imaginary argument of the 1-th order; $x = q/v(v^2t^2 - 3r^2)^{1/2}$. The expressions (19), (20) determine a_0 and a_1 at $v^2t^2 > 3r^2$; $a_0 = a_1 = 0$ at $v^2t^2 < 3r^2$. For great t formulas (19), (20) pass asymptotically to formulas of diffusive approximation. Graphs of a_0 and a_1 dependence on r at $qt = 3$ (solid curves) are plotted in Fig.1. For comparison we brought out also the values of a_0 and a_1 in diffusive approximation (dashed curves). As may be seen, the main discrepancy is obtained in fluxes.

Inasmuch as in the initial equations we neglected the terms containing a_2, the condition of (19), (20) applicability will be the smallness of a_2. This quantity may be determined by an approximate formula obtained from the third equation of the system (10)

$$a_2 \approx -\frac{1}{9}\frac{v}{q}\frac{\partial a_1}{\partial r} + \frac{1}{9}\frac{v}{qr}a_1.$$

Substituting its value according to formula (20) in place of a_1 we shall have

$$a_2 = -\frac{3\sqrt{3}\bar{N}q^5r^2}{16\pi^2 v^5 x^3}e^{-qt}I_1(x)\left(1 - \frac{4I_0}{I_1 x} + \frac{8}{x^2}\right). \tag{21}$$

The estimate of a_2 shows that (19), (20) are valid through $r = vt/\sqrt{3}$ at $qt \leqslant 5$. For great t, these expressions are valid only for $r < r_0 < vt/\sqrt{3}$, at the same time, this inequality is strenghtened with the increase of t. The distribution of particles at $vt > r > vt/\sqrt{3}$ is given by the δ- function. This is the corollary of the unaccounted a_2. Therefore, for the description of particles found at $r > r_0$, it is necessary to take onto account a_2, a_3 and so forth., that is, in the zone $r > r_0$ the particles are distributed anisotropically.

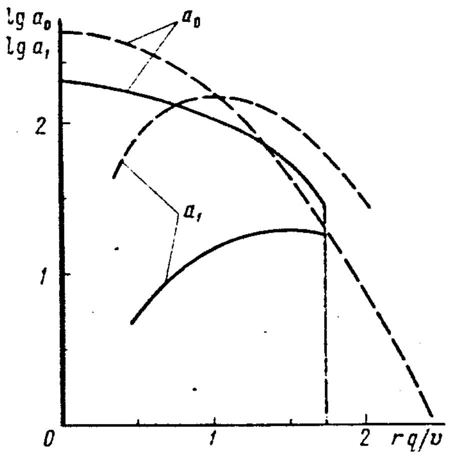

Fig. 1

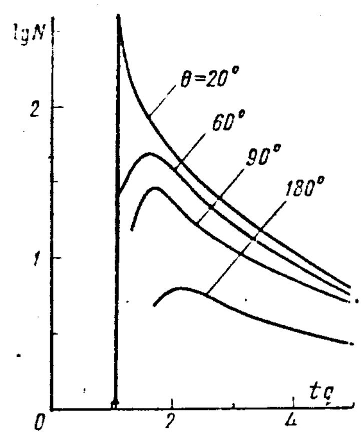

Fig. 2

Let us pass now to the question of angular distribution of particles in the zone $r > r_0$. For the solution of this problem we shall substitute the derivative in respect to θ by finite differences, using the method of straight lines [2]. As a result, we shall obtain instead of the equation (6) a system of equations in partical derivative of first order

$$\frac{\partial N_i}{\partial t} + v\cos\theta_i \frac{\partial N_i}{\partial r} - \frac{v\sin\theta_i}{2rh}(N_{i+1} - N_{i-1}) =$$
$$= \frac{1}{h^2}(N_{i+1} - 2N_i + N_{i-1}) + \frac{1}{2h}\operatorname{ctg}\theta_i (N_{i+1} - N_{i-1}). \tag{22}$$

Here h is a step; $N_i \doteq N(r, \theta_i, t)$; $i = 0, 1, 2, \ldots$; N_i is the mean value of the function $N(r, \theta, t)$, averaged in the interval $\theta_i - {}^1/_2 h, \theta_i + {}^1/_2 h$.

Subsequently, we shall make the following approximations: inasmuch as in the considered zone the particles are strongly anisotropically distributed, we shall neglect in the equation (22) the terms containing N_2; since the thickness of the considered zone is less than the distance to the center when t is not too great, we shall also disregard the sphericity.

On the whole, we have the following system:

$$\frac{\partial N_0}{\partial t} + v0{,}94\frac{\partial N_0}{\partial r} = -q4{,}02N_0 + q4{,}02N_1,$$
$$\frac{\partial N_1}{\partial t} + v0{,}50\frac{\partial N_1}{\partial r} = q1{,}64N_0 - q4{,}10N_1. \tag{23}$$

Here the coefficients are computed at $h = 40^{\circ}$, $\theta_0 = 20^{\circ}$. If for initial conditions we take

$$N_0|_{t=0} = \frac{\overline{N}}{16\pi^2 r^2}\delta(r),\ N_1|_{t=0} = 0, \tag{24}$$

the solution of the system has the form

$$N_0 = \frac{4{,}54}{16\pi^2 r^2}\delta(t-\zeta)e^{-2{,}57qt} +$$
$$+\frac{5.84q\overline{N}}{16\pi^2 v^2 r^2}\frac{t+\zeta}{\sqrt{t^2-\zeta^2}}I_1(2{,}57q\sqrt{t^2-\zeta^2})\exp[-4{,}06qt - 0{,}04q\zeta],$$
$$N_1 = \frac{3{,}71\overline{N}q}{16\pi^2 v r^2}I_0(2{,}57q\sqrt{t^2-\zeta^2})\exp[-4{,}06qt - 0{,}04\zeta] \tag{25}$$
$$\xi = 4{,}54\left(\frac{r}{v} - 0{,}72t\right),$$

where $N_0, N_1 = 0$ at $t^2 < \zeta^2$. This solution describes satisfactorily the distribution of particles at $qt < 1$, when the velocities of the principal mass of particles are directed forward. At great t it is applicable only for the outermost parts of the zone $r > r_0$. For the inner part of this zone it is necessary to account for the particles having drifted into the zone $r < r_0$ and having caught up with the zone $z > z_0$ again after a series of scatterings. In order to take these particles into account, we shall proceed as follows. We shall choose a distance r_1, such that formulas (19) and (20) be still valid, and that at the same time there be a sufficiently strong anisotropy in the angular distribution of particles (that is, a_0 and a_1 are comparable). For not too great t we may choose for r_1 $vt/\sqrt{3}$. Having expressed N_0 and N_1 by a_0 and a_1 at the point $r = r_1$, we shall obtain the boundary conditions for the system (23):

$$N_0|_{r=vt/\sqrt{3}} = (a_0 + 0{,}94\,a_1)|_{r=vt/\sqrt{3}} = \frac{1}{0{,}36\,v}p_0(t),$$
$$N_1|_{r=vt/\sqrt{3}} = (a_0 + 0{,}50\,a_1)|_{r=vt/\sqrt{3}} = \frac{1}{0{,}08v}p_1(t). \tag{26}$$

The quantities p_i have the sense of power of kind-i particle sources, moving with the velocity $v/\sqrt{3}$. Resolving the system (23) with zero initial conditions and boundary conditions (26), we shall obtain

$$N_0 = \frac{q}{v}\int_0^t p_0(\tau)\frac{t^2}{1{,}5\tau^2}\overline{N}_0\left(r-\frac{v\tau}{\sqrt{3}},\ t-\tau\right)d\tau,$$

$$N_1 = \frac{q}{v}\int_0^t p_0(\tau)\frac{t^2}{1{,}5\tau^2}\overline{N}_1\left(r-\frac{v\tau}{\sqrt{3}},\ t-\tau\right)d\tau. \tag{27}$$

Here N_0 and N_1 are determined by formulas (25), the multiplier r^{-2} being rejected. The multiplier $t^2/1.5\tau^2$ approximately accounts for the sphericity. The results of numerical calculations of time dependence of particle intensity in the directions $20, 60, 100, 180^o$ over distances $rq/v = 1, 3, 5$ are plotted in Figures 2 – 4. It may be seen from the graphs that for a spherically-symmetrical flare strong anisotropy in the angular distribution of particles will be observed only at $rq/v \leqslant 1$ at time of onset of particle density maximum (that is, N_0 and N_1 differ substantially).

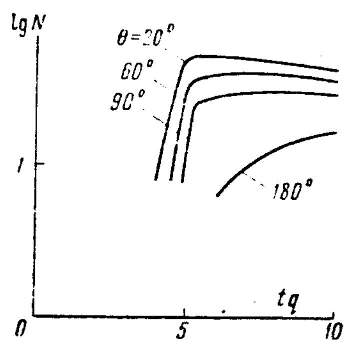

Fig. 3

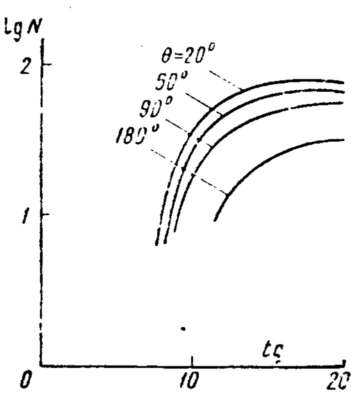

Fig. 4

<u>AXISYMMETRICAL FLARE</u>. – Let us consider a flare with a directed ejection. We shall assume that particles are ejected at the point $r = 0$ along tha axis $\eta = 0$. In this case the kinetic equation has the following form

$$\frac{\partial N}{\partial t} + v\cos\theta\frac{\partial N}{\partial r} - v\frac{\sin\theta}{r}\frac{\partial N}{\partial\theta} + v\frac{\sin\theta\cos\varphi}{r}\frac{\partial N}{\partial\eta} -$$

$$- v\frac{\sin\theta\sin\varphi}{r}\operatorname{ctg}\eta\frac{\partial N}{\partial\varphi} = \frac{q}{\sin\theta}\frac{\partial}{\partial\theta}\sin\theta\frac{\partial N}{\partial\theta} + \frac{q}{\sin^2\theta}\frac{\partial^2 N}{\partial\varphi^2} + \omega\frac{\partial N}{\partial\varphi}. \tag{28}$$

Representing N in the form of series

$$N = b_0(r,\eta,\theta,t) + \sum_{k=1}^{\infty}[b_k(r,\eta,\theta,t)\cos k\varphi + c_k(r,\eta,\theta,t)\sin k\varphi] \tag{29}$$

we shall obtain in the diffusive approximation for the propagation along the coordinate η the equation determining b_0

$$\frac{\partial b_0}{\partial t} + v\cos\theta\frac{\partial b_0}{\partial r} - \frac{v\sin\theta}{r}\frac{\partial b_0}{\partial \theta} =$$
$$= \frac{qv^2\sin^4\theta}{4r^2(\omega^2\sin^4\theta + q^2)}\frac{1}{\sin\eta}\frac{\partial}{\partial\eta}\sin\eta\frac{\partial b_0}{\partial\eta} + \frac{q}{\sin\theta}\frac{\partial}{\partial\theta}\sin\theta\frac{\partial b_0}{\partial\theta}. \tag{30}$$

The following initial condition must be added to the equation (30):

$$b_0|_{t=0} = \frac{\bar{N}}{16\pi^2r^2} = \delta(r)\delta(1-\cos\eta)\delta(1-\cos\theta), \tag{31}$$

alongside with the condition of periodicity with respect to η and θ.

Estimating that $\omega \gg q$, we see that the coefficient at the derivative with respect to η does not depend on θ, provided we estimate also that it is independent from r (that is, taking the effective average), the variable η is determined from the remaining in the form

$$b_0 = N_{c\varphi}\frac{1}{2\bar{D}_\eta t}\exp\left(-\frac{\eta^2}{4\bar{D}_\eta t}\right). \tag{32}$$

Here $N_{c\varphi}$ is the solution for a spherically-symmetrical flare; $\bar{D}_\eta$ is the effective value of the coefficient at the derivative with respect to

$$\bar{D}_\eta = \overline{qv^2}/4r^2\omega^2.$$

We shall consider the question as to how the flare's nonsphericity can manifest itself upon the observation condition of strong anisotropy. As may be seen from the earlier presented graphs, in the case of spherically-symmetric flare the intensity rises extremely rapidly in the anisotropic part of the curve and the multiplier $1/\bar{D}_\eta t$ can not influence somewhat substantially the shape of the curve in that part. In Fig. 5 we brought out a graph corroborating this remark. It indicates the course of intensity with time for $rq/v = 3$, $\eta = 0$. The exponential multiplier $\exp(-\eta^2 > 4\bar{D}_\eta t)$ may shift the maximum to the isotropic region in the case when $\eta^2 > 4\bar{D}_\eta t$. Therefore, the strong anisotropy will be observed only in the case when the observed is located inside the cone $rq/v \leqslant 1$, $\eta^2 < 4\bar{D}_\eta t$.

COMPARISON WITH OBSERVATIONS.- Let us compare the conclusions obtained with the observation data on the flare of 4 May 1960 brought out in [3]. The symmetry axis direction of the flux of arriving protons deviated by 50^o to the west from the direction at the Sun. Apparently the particles were led by the lines of force of a regular spiral-like field; this is why in the given case the angle must be counted from the direction of the line of force of the regular field at the given point and not from the radial direction, as the former coincides with the symmetry axis of the flux. Taking this remark into account we shall apply the results obtained for the radial field to the real, estimating that other variations are immaterial. According to data of [3], a strong anisotropy was preserved in the course of a half-hour after onset of particle density maximum. This means, first of all, that at time of that flare, the Earth was located in the favorable cone. Secondly, the duration of the strong anisotropy was determined by the duration of source's operation, for, as may be seen from the above graphs, the duration of strong anisotropy does not exceed a few minutes in an instantaneous source. Plotted in Fig. 6 are the values of intensity in the directions of 40 and 60^o as a function of time at $q = 2 \cdot 10^{-3} s^{-1}$, $r = 1.5 \cdot 10^{13}$ cm, $\eta = 0$. The power of the source had an exponential dependence on time $\exp(-t/t_o)$, where $t_o = 16$ min. The values of fluxes observed at the respective points on Earth are plotted by circles and triangles. It may be seen that the curve brought out describes satisfactorily the force and the duration of the anisotropy.

t_o may denote the source's operation time as well as the deexcitation time of the region of diffusion in Sun's vicinity. Let us pause at the second variant. The deexcitation time of a sphere of radius R, filled with a medium with diffusion coefficient D_R, is approximately equal to R^2/D_R. Assuming $R = 10R_{\odot} = 7 \cdot 10^{11}$cm and estimating the deexcitation time to be equal to 16 min., we shall obtain $D_R = 5 \cdot 10^{20}$ $cm^2 \cdot sec$. For the farther parts of the interplanetary space the diffusion coefficient is equal to $D_A = 5 \cdot 10^{22}$ $cm^2 sec$. Obviously it is not indispensable to estimate that the near-solar diffusion region is sharply outlined. It is more plausible to postulate a smooth variation of the diffusion coefficient as the distance from the Sun varies. In this case D_R is a certain effective value of the diffusion coefficient at small distances from the Sun, whereas D_A is its mean value at more remote distances, through the Earth's orbit. Note that inasmuch as the anisotropy region was investigated by a curve, the medium beyond the Earth's orbit did not practically affect the values of the parameters being determined. The diffusion coefficient, determined over the dropping part of the curve, refers mainly to the space beyond the Earth's orbit. According to the data of [4], it is equal to $3 \cdot 10^{22}$ cm^2.sec. Therefore, at $r < 1$ a.u. the diffusion coefficient rises rapidly with the distance (approximately as r^2), while at $r > 1$ a.u. it is nearly constant. The last circumstance may be explained by the influence of the regular field, which beyond the Earth's orbit becomes essentially azimuthal.

Considering the value of the mean square of irregular field's intensity as known, we may be able to determine the cha-

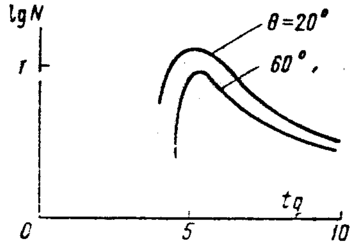

Fig. 5

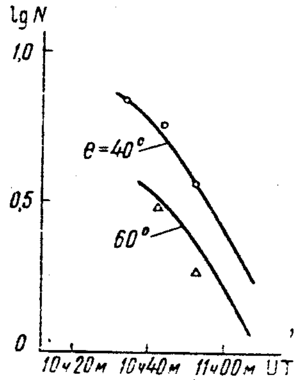

Fig. 6

racteristic dimension of the irregularities over which the relativistic protons are scattered. Assuming $\sqrt{b^2} = 3 \cdot 10^{-5}$ gauss, we obtain $a = 3 \cdot 10^{10}$ cm.

The fact that other flares do not indicate such a strong anisotropy, even those which had heliocoordinates nearly identical to those of the flare of 4 May 1960, implies that the span of the favorable cone does not exceed 10 - 20°.

In conclusion we shall formulate the propagation pattern of high-energy solar protons. Near the Sun, the dimension of irregularities are apparently greater than the Larmor radius, and the propagation process is represented by random drifts; here, the diffusive approximation may be applied. At distances of several tens of solar radii the effective free path becomes comparable with the distance to the center and this is why the diffusion approximation is here inapplicable; the Larmor radius becomes greater than the characteristic dimension of irregularities and particles scatter over a small angle. Inasmuch as the effective free path remains constant, the diffusion regime settles again during a time of the order of the time of the free path.

**** THE END ****

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NASA HQS

SS	NEWELL, NAUGLE
SG	MITCHELL
	ROMAN
	SMITH
	SCHARDT
	DUBIN
SL	LIDDEL
	FELLOWS
	HIPSHER
	HOROWITZ
SM	FOSTER
	GILL
RR	KURZWEG
RTR	NEILL
ATSS - T	
WX	SWEET

OTHER CENTERS

AMES R C

SONETT
LIBRARY

LANGLEY R C

160	ADAMSON
116	KATZOFF
185	WEATHERWAX

JPL

SNYDER
NEUGEBAUER
WYCKOFF

UNIV. IOWA

VAN ALLEN

U C BERKELEY

WILCOX

U C L A

COLEMAN